

Proposal for a Constant Cosmological Constant

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Abstract

It is proposed that the apparent positive acceleration of the cosmological scale factor is due to the vacuum energy of an incomplete chiral phase transition in a hidden $SU(2)$ sector. Constraints from primordial nucleosynthesis imply that the present metastable phase is in a substantially supercooled state. It is argued that massless chiral condensates can substantially enhance the possibility of supercooling, and a linear sigma model exhibiting scale invariance broken only at the quantum level is shown to accommodate the required supercooling with a reasonable choice of quartic couplings. The extensive supercooling can in principle be confirmed or rejected on the basis of interface tension measurements in lattice simulations with dynamical fermions.

1 Introduction

There is accumulating evidence that the expansion rate of the universe is greater now than in the past [1]. The data, taken alone or in conjunction with constraints from large scale structure [2], are compatible (in a flat universe) with a contribution $\Omega_\Lambda \simeq 0.7$ from a cosmological constant Λ or its equivalent. Cluster abundance estimates for the matter fraction Ω_{matter} of the critical density, combined with an analysis [3] of CMB observations in the Doppler peak region (which support the inflation prediction $\Omega_{tot} = 1$) are also compatible with such a contribution to the “dark energy”. The resulting energy density is given by

$$\rho_\Lambda = (2.2 \times 10^{-3} \text{ eV})^4 \cdot (\Omega_\Lambda/0.7) \cdot (h/0.65)^2 \quad , \quad (1)$$

where h = present Hubble constant H_0 in units of 100 km/sec/Mpc. The introduction of a non-zero value for ρ_Λ or its equivalent presents an important challenge to particle physicists and cosmologists. Even if one concedes ignorance and simply accepts as premise that the universe is relaxing to a state with $\rho_\Lambda = 0$ [4], there still remains the perplexing question of the origin of a mass scale $\rho_\Lambda^{1/4} \sim 10^{-3} \text{ eV}$ which makes ρ_Λ relevant in the present era. Discussion in recent years has centered on models in which Λ becomes time-dependent, originating in the energy density of a scalar field ϕ evolving in a potential $V(\phi)$. One candidate for ϕ is a pseudo-Nambu-Goldstone boson (axion) in a harmonic potential associated with the breaking of a $U(1)$ symmetry to Z_N [5]. The cosmological consequences of this scenario depend on both the normalization of the potential M^4 and the scale f at which the $U(1)$ symmetry is realized in the Goldstone mode, and there have been studies [6, 7] where the parameters make the model relevant to late-time cosmology, including recent applications to the SNeIa results [8, 9]. The required scale $M \sim 10^{-3} \text{ eV}$ can be associated with a neutrino mass [6] or with the confining scale of a hidden gauge theory [10], ϕ being the axion. In quintessence models [11, 12], ϕ evolves so that the late-time behavior of the dark energy density ρ_ϕ is largely independent of initial conditions. What remains to be tuned by hand is the parameter in the potential which allows for a positive acceleration of the scale parameter during the SNeIa era relevant to the observations of Ref. [1]. The origin of the field ϕ , the form of its potential, and its relation to other physics remain to be explained [13].

In this paper, I would like to propose that the dark energy ρ_Λ is not evolving, but is the false vacuum energy associated with an incomplete chiral phase transition in a hidden $SU(2)'$ gauge theory with strong scale $\sim \rho_\Lambda^{1/4}$. It will be seen that in the context of modern D-brane physics, the scale $\rho_\Lambda^{1/4}$ for the vacuum energy can be accommodated in a natural manner in a supersymmetric GUT theory. Instead, the central problem will be to explain how a low temperature $T_{hidden} < T_{CMB} \simeq \frac{1}{10}\rho_\Lambda^{1/4}$ can be sustained for the quark-gluon phase of the (supercooled) plasma of the hidden sector. Such a high degree of supercooling can potentially be tested in lattice simulations. In the present work, it will strongly constrain the effective field theory. The model is described in the next section, and some possible advantages as an alternative to the scalar field scenario are mentioned in the concluding section.

2 The Model

I will consider a hidden unbroken $SU(2)'$ Yang-Mills theory whose low energy matter content is N_f Dirac fermions ($2N_f$ Weyl fermions) with vector-like coupling to the gauge field. (For now, hidden sector quantities will be denoted by a prime.) Except for gravitational interactions, this $SU(2)'$ is completely decoupled from all standard model fields. The choice of $SU(2)'$ is doubly motivated: (1) the running of the gauge coupling is slow, so that the scale $\Lambda_{SU(2)'}$ can be pushed to values approximating $\rho_\Lambda^{1/4}$ [6] (2) there will be fewer massless degrees of freedom to perturb the successful scenario of Big Bang Nucleosynthesis (BBN), a critical requirement of the model. The hidden nature of the $SU(2)'$ will be discussed when the coupling requirement at GUT energies is obtained.

Since the theory will be considered to have evolved from GUT energies, the supersymmetric version becomes relevant. The matter content then consists solely of $2N_f$ chiral $SU(2)'$ doublet superfields $Q_i (i = 1 \dots 2N_f)$. In order to preserve the low energy chiral phase transition, $SU(2)'$ singlet, flavor antisymmetric mass terms $\sim Q_i Q_j - Q_j Q_i$ must be prohibited by a discrete symmetry (R -invariance or $Z_N, N > 2$ symmetry). Next, the 1-loop RG equation relates the gauge coupling at GUT α'_{GUT} and the strong coupling scale $\Lambda_{SU(2)'}:$

$$\begin{aligned} \frac{2\pi}{\alpha'_{\text{GUT}}} &= b_2^{\text{SUSY}} \ln \left(\frac{M_{\text{GUT}}}{1 \text{ TeV}} \right) + b_2^{\text{non-SUSY}} \ln \left(\frac{1 \text{ TeV}}{\Lambda_{SU(2)'}} \right) \\ b_2^{\text{SUSY}} &= 6 - N_f \\ b_2^{\text{non-SUSY}} &= \frac{22}{3} - \frac{2}{3} N_f \ , \end{aligned} \quad (2)$$

so that for $M_{\text{GUT}} = 2 \times 10^{16} \text{ GeV}$, $\Lambda_{SU(2)'} = \rho_\Lambda^{1/4} = 2.2 \times 10^{-3} \text{ eV}$ one obtains

$$\alpha'_{\text{GUT}}{}^{-1} = 68.6 - 8.46 N_f \ . \quad (3)$$

The major premise of the model is the existence of a false vacuum at present. A first order phase transition driven by fluctuations of chiral condensates is strongly indicated by theoretical arguments [14, 15, 16], most cleanly for $N_f \geq 4$. From Eq. (3), it would seem that unification with standard model couplings ($\alpha_{\text{GUT}} \simeq \frac{1}{25}$) can be achieved with a choice $N_f = 5$; however, this would elevate the number of Goldstone degrees of freedom below the critical temperature ($= 2N_f^2 - N_f - 1$) [18] to a value larger than the effective number of quark-gluon degrees of freedom ($= 7N_f + 6$), presumably vitiating the first order phase transition. Thus I choose $N_f = 4$, and Eq. (3) gives

$$\alpha'_{\text{GUT}} \simeq \frac{1}{35} \neq \alpha_{\text{GUT}} \ . \quad (4)$$

Such a disparity in GUT-scale gauge couplings is not difficult to accommodate in current formulations of string theory, with gauge fields residing in open strings tied to D-branes. If, for example, the standard model gauge group lives on a 5-brane and the hidden $SU(2)'$ on another 5-brane (orthogonal with respect to the compactified 2-tori) [19], the ratio of the gauge couplings would be inversely proportional to the volumes of the 2-tori:

$$\frac{\alpha'_{\text{GUT}}}{\alpha_{\text{GUT}}} = \frac{v_2}{v'_2} \ . \quad (5)$$

Thus, a 20% difference in toroidal moduli could account for the disparity in the α 's.

Temperature Constraints

Having established a model, one can quickly ascertain the constraints which follow from BBN. The hidden sector energy density ρ' of the $SU(2)'$ gauge fields and $2N_f$ Weyl doublets, relative to a single species of left-handed neutrino is given by

$$\left. \frac{\rho'}{\rho_{\nu_e}} \right|_{BBN} = \left(\frac{7N_f + 6}{(7/4)} \right) \left(\frac{T'}{T} \right)^4 \Big|_{BBN} . \quad (6)$$

Requiring this ratio to be ≤ 0.3 [20] implies (for $N_f = 4$)

$$\left. \frac{T'}{T} \right|_{BBN} \leq 0.353 . \quad (7)$$

Much of this can be accounted for through reheat processes in the visible sector. Assuming no reheat for the hidden sector for energies below the electroweak scale, one finds

$$\left. \frac{T'}{T} \right|_{BBN} = \left(\frac{g^*(BBN)}{g^*(>EW)} \right)^{1/3} \left. \frac{T'}{T} \right|_{>EW} \quad (8)$$

$$= \left(\frac{10.75}{106.75} \right)^{1/3} \left. \frac{T'}{T} \right|_{>EW} = 0.465 \left. \frac{T'}{T} \right|_{>EW} , \quad (9)$$

where $>EW$ denotes temperatures above the electroweak scale. In order to comply with the BBN requirement (Eq. 6) an additional suppression of T'/T above the electroweak scale by a factor of 0.76 is required. Here one can invoke an asymmetric post-inflation reheat into visible and hidden sector quanta [21]. In the slow reheat scenario [22], T_{reheat} is proportional to the coupling of the inflaton to the quanta [23], so that a ratio of couplings of the same order as the ratio of the α 's (4) could provide the desired additional suppression. This reheat asymmetry could have the same origin as the α asymmetry, if the inflaton originates in the modular sector. Asymmetric reheating also obtains in the case of parametric resonance [24] because of the asymmetric coupling [21].

Because of e^+e^- annihilation, the present T' is further depressed relative to the present CMB temperature by the same $(4/11)^{1/3}$ factor as with neutrinos. Thus, together with (7), one obtains $(T'_{\text{now}}/T_{\text{CMB}}) \leq 0.251$. Since $T_{\text{CMB}} = 2.35 \times 10^{-4} \text{ eV} = 0.11 \rho_\Lambda^{1/4}$, one finds

$$T'_{\text{now}}/\rho_\Lambda^{1/4} \leq 0.028 . \quad (10)$$

This suggests a great deal of supercooling, which needs to be accommodated. The calculations in this work will require the ratio T/T_c (T_c is the critical temperature¹) which in turn requires knowledge of the ratio $\rho_\Lambda^{1/4}/T_c$. This will be calculable in the effective field theory to be discussed.

Supercooling

In the standard formulation of first order phase transitions via critical bubble formation [25, 26] the condition for *failure* to complete a phase transition in the expanding universe at (hidden)

¹From here on, all temperatures will be understood to be hidden sector temperatures, and the primes will be omitted.

temperature T is [27]

$$(T/H(T))^4 e^{-F_c/T} < 1 \quad , \quad (11)$$

where F_c is the free energy of a critical bubble, and $H(T)$ is the Hubble constant at temperature T . With $H_0 \simeq 2.2 \times 10^{-33} \text{ h eV}$, $T = 0.28 T_{\text{CMB}}$, one obtains the condition for failure to nucleate

$$F_c/T > 260 \quad . \quad (12)$$

In the thin wall approximation, the bubble has a well-defined surface tension σ , and the picture is consistent only for small supercooling below T_c . The bubble action is given by [25, 26]

$$\frac{F_c}{T} = \frac{16\pi}{3} \frac{\sigma^3}{L^2 \eta^2 T_c} \quad , \quad (13)$$

where L is the latent heat and $\eta = (T_c - T)/T_c$. Thus, a failure to nucleate via thin-walled bubbles requires a large surface tension, $\sigma/T_c^3 \gtrsim 1$. In lattice studies of quenched QCD [28], the interface tension between confined and deconfined phases is small: $\sigma/T_c^3 \simeq 0.1$. However, a simple calculation [29] based on the MIT bag model [30] suggests that the picture can change drastically in the presence of chiral condensates: in that case,

$$\sigma = -\frac{1}{4} \sum_{i=1}^{N_f} \langle \bar{q}_i q_i \rangle \quad . \quad (14)$$

which for QCD ($T_c \simeq 150 \text{ MeV}$, $\langle \bar{q}_i q_i \rangle \simeq -(240 \text{ MeV})^3$ per flavor) would give a large surface tension, $\sigma \simeq 4T_c^3$, possibly invalidating the thin-wall approximation.² A full analysis of the unquenched QCD situation is probably best carried in the framework of a mean field theory [33]. I will proceed in the context of such a theory to see what constraints are imposed on the $SU(2)'$ model in order to attain the desired metastability (Eq. (12)) until the present era.

3 Linear Sigma Model for $SU(2)'$ with N_f Flavors.

The symmetry breaking pattern of color $SU(2)'$ with N_f flavors has long been known [18], and a linear sigma model for this case has recently been examined [34]. Such a model will serve conveniently to study the chiral phase transition. The meson and diquark baryon fields are contained in the $2N_f \times 2N_f$ antisymmetric matrix $\Phi_{ij} = -\Phi_{ji}$, with chiral symmetry breaking occurring in $Sp(2N_f)$ direction [18] compatible with the Vafa-Witten theorem [35, 18, 36]

$$\langle \Phi \rangle_0 = \frac{\phi_0}{\sqrt{2(2N_f)}} \begin{pmatrix} \mathbf{0} & \mathbf{1} \\ -\mathbf{1} & \mathbf{0} \end{pmatrix} \quad (15)$$

where $\mathbf{0}$ and $\mathbf{1}$ are $N_f \times N_f$ matrices. The lagrangian is

$$\mathcal{L} = \text{Tr } \partial^\mu \Phi \partial_\mu \Phi - m^2 \text{Tr } \Phi^\dagger \Phi - \lambda_1 (\text{Tr } \Phi^\dagger \Phi)^2 - \lambda_2 \text{Tr } \Phi^\dagger \Phi \Phi^\dagger \Phi \quad . \quad (16)$$

I have omitted a term $\propto \text{Pf}(\Phi) + \text{Pf}(\Phi^\dagger)$ arising from the axial anomaly. Since I will be working with $N_f = 4$, this additional operator quartic in the fields will not qualitatively change the discussion which follows. Stability in all field directions requires $\lambda_2 \geq 0$, $\lambda_1 + \lambda_2/2N_f \geq 0$.

²This would give an average distance between nucleation sites of about 10 m [31], perhaps marginally affecting primordial light element abundances [32].

Specializing now to the field ϕ in the direction of the vev, one obtains

$$\begin{aligned}\mathcal{L}_0 &= \frac{1}{2}(\partial_\mu\phi)^2 - V(\phi) \ , \\ V(\phi) &= \frac{1}{2} m^2\phi^2 + \frac{1}{4} \bar{\lambda}\phi^4 \ , \\ \bar{\lambda} &= \lambda_1 + \bar{\lambda}_2, \quad \bar{\lambda}_2 = \lambda_2/2N_f \ .\end{aligned}\tag{17}$$

With finite temperature corrections (restricted for simplicity to the m^2 term) and the introduction of a running quartic coupling, one obtains the effective potential

$$V(\phi, T) = \frac{1}{2} m^2(T)\phi^2 + \frac{1}{4} \bar{\lambda}(t)\phi^4 \ ,\tag{18}$$

where $t = \ln(\phi/\phi_0)$, and m^2 has the standard T -corrected form

$$m^2(T) = A(T^2 - T_0^2) \ .\tag{19}$$

In the model described, A can be calculated, and I find at one loop

$$A = \frac{1}{12} [(2N_f(2N_f - 1) + 2)\bar{\lambda} + (6N_f(2N_f - 1) - 2)\bar{\lambda}_2] \ .\tag{20}$$

The ϵ expansion analysis of the model described by (16) shows that it allows a first order phase transition through a Coleman-Weinberg mechanism at $T = T_0$, when $m^2(T) = 0$ [34]. However, it will shortly be apparent that the large supercooling will require that $T_0^2 \ll T_c^2$. This (approximate) conformal invariance at tree level in the chiral lagrangian (to be discussed below) in turn suggests that chiral symmetry breaking at *zero temperature* in this model also proceeds through radiative corrections (Coleman-Weinberg) [17]. Thus, to lowest order (see Eq. (29) below)

$$\bar{\lambda}(t) = -\lambda(1 - 4t) \ , \quad \lambda \equiv -\bar{\lambda}(0) \ .\tag{21}$$

where $t = 0$ is defined by the minimum of the second term in (18).

The vacuum at $\phi = 0$ described by the potential (18) becomes metastable at a temperature T_c determined by requiring simultaneously

$$V'(\phi_+) = V(\phi_+) = 0 \ ,\tag{22}$$

at some field value ϕ_+ . A short algebraic exercise with Eqs. (18), (19), (21), and (22) determines T_c :

$$m(T_c) = \sqrt{A(T_c^2 - T_0^2)} = \sqrt{\lambda}\phi_0 e^{-1/4} \ .\tag{23}$$

Suppression of m^2 .

For $T_0 < T < T_c$ there is a barrier between the false and true vacua. But the extreme supercooling requirement indicated in Eq. (10) will be seen to impose a large hierarchy, $T_0 \ll T_c$ (it will turn out that $T_0 \leq 0.124T_c$.) There is no obvious argument to justify this hierarchy — it is simply required in this model for compatibility with the supercooling requirement. Nevertheless, two comments may be made: (1) A quantitative comparison with the linear $SU(3) \times SU(3)$ sigma model (whose

dynamics differs through the presence of the cubic term) is perhaps instructive. In that case, the coefficient A is obtained as a sum over the massive mesons [37], $A = \frac{1}{12} \sum_i M_i^2/v^2$, where $v \simeq f_\pi$. With $M_i \simeq 1$ GeV, one estimates $A \simeq 83$. Since $T_0 = \sqrt{-\mu^2/A}$, where $\mu^2 \simeq -0.15$ GeV² [38] is the temperature-independent (negative) mass parameter, we find $T_0 \simeq 43$ MeV $\simeq 0.24T_c$. This may provide a normative expectation for T_0/T_c . (2) In $SU(N)$ gauge theories with N_f massless fermion flavors, a *continuous* transition to an approximately conformal phase (which would imply $m^2 = 0$) is suggested at some value of $N_f/N < 11/2$ [39]. Some analytic studies [39, 40] indicate $N_f/N \approx 4$ as a critical value, while a QCD lattice study [41] show hints of a suppression of the chiral condensate (expected in the transition to the conformal phase [39]) for a smaller value, $N_f/N = \frac{4}{3}$ ($SU(3)$ with 4 flavors). Perhaps the ratio in the present case ($N_f/N = 2$) is sufficiently large to significantly suppress the zero temperature m^2 term in the effective lagrangian — a lattice study of chiral symmetry breaking in $SU(2)$ with $N_f = 4$ could in principle shed light on this question. At any rate, at this juncture I accept the hierarchy $T_0 \ll T_c$, and simplify matters even more by setting

$$T_0 = 0 \quad (24)$$

in Eq. (19). In such a model, with *two* coupling constants, the chiral invariance is broken at $T = 0$ in the Coleman-Weinberg manner [17, 42]. For $T \neq 0$, the transition becomes first order, as described in the previous section.

Critical Bubbles.

For $T \neq 0$ it will prove convenient to rescale $\phi = m(T)\phi'/2\sqrt{\lambda}$, so that using Eqs. (18), (21) and (23) one may write

$$\begin{aligned} V &\equiv \frac{m^4(T)}{4\lambda} \bar{V} \\ \bar{V} &= \frac{1}{2}\phi'^2 + \frac{1}{4}\phi'^4 \left[\ln((m(T)/m(T_c))\phi')/2 - \frac{1}{2} \right] \\ &= \frac{1}{2}\phi'^2 + \frac{1}{4}\phi'^4 \left[\ln((T/T_c)\phi'/2) - \frac{1}{2} \right] \end{aligned} \quad (25)$$

for $T_0 = 0$.

The O(3) symmetric free energy for a critical bubble formed at temperature T is given by [43]

$$\begin{aligned} F_c &= 4\pi \frac{m(T)}{4\lambda} \int_0^\infty dr' r'^2 \left[\frac{1}{2} (d\phi'/dr')^2 + \bar{V}(\phi', T/T_c) \right] \\ &\equiv \frac{m(T)}{4\lambda} f(T/T_c) \end{aligned} \quad (26)$$

where $r' = m(T)r$. The field ϕ' is the solution to

$$\frac{d^2\phi'}{dr'^2} + \frac{2}{r'} \frac{d\phi'}{dr'} = \frac{\partial \bar{V}}{\partial \phi'} \quad (27)$$

subject to $d\phi'/dr'|_{r'=0} = 0$, $\phi'(\infty) = 0$. With the help of Eqs. (19), (20) and (24) the bubble action can then be calculated more explicitly in terms of the the quantity $f(T/T_c)$. For $N_f = 4$, I obtain

$$\begin{aligned}
F_c/T &= \sqrt{A}/(4\lambda) f(T/T_c) \\
&= \sqrt{\frac{29}{6}\bar{\lambda} + \frac{83}{6}\bar{\lambda}_2}/(4\lambda) f(T/T_c) .
\end{aligned} \tag{28}$$

At this point, I implement the condition of radiative symmetry breaking (at $T = 0$) in the two-parameter $(\bar{\lambda}, \bar{\lambda}_2)$ space. This imposes the condition [44]

$$-4\bar{\lambda}(0) = 4\lambda = \beta_{\bar{\lambda}}(0) . \tag{29}$$

From the one-loop RG equations [34], for $\bar{\lambda}_2(0) \gg |\bar{\lambda}(0)|$, but $\bar{\lambda}_2(0)$ still perturbative (this will be justified *a posteriori*), Eq. (29) gives

$$\bar{\lambda}_2(0) = \left(4\pi/\sqrt{4N_f^2 - 2N_f - 2}\right) \sqrt{-\bar{\lambda}(0)} = (4\pi/\sqrt{54})\sqrt{\bar{\lambda}} , \tag{30}$$

and hence in the same approximation

$$A \simeq \frac{83}{6} \frac{4\pi}{\sqrt{54}} \sqrt{\bar{\lambda}} . \tag{31}$$

Thus the bubble action can be written entirely in terms of the coupling constant $\lambda = -\bar{\lambda}(0)$:

$$\begin{aligned}
F_c/T &= \sqrt{\frac{4\pi}{\sqrt{54}} \frac{83}{6} \sqrt{\bar{\lambda}}}/(4\lambda) f(T/T_c) \\
&\simeq 1.22 f(T/T_c) \lambda^{-3/4} .
\end{aligned} \tag{32}$$

We now require the ratio T/T_c . From Eq. (10), one needs to relate the vacuum energy ρ_Λ to T_c . At $T = 0$, $m(T) = 0$, Eqs. (18), (21), (23), (24) and (31) give

$$\begin{aligned}
\rho_\Lambda &= V(0) - V(\phi_0) \\
&= \frac{1}{4}\lambda \phi_0^4 \\
&= \frac{1}{4}eA^2T_c^4/\lambda \\
&= (4.4T_c)^4 .
\end{aligned} \tag{33}$$

Combining this with Eq. (10), we have the supercooling requirement

$$T/T_c \leq 0.124 . \tag{34}$$

The bubble action may now be evaluated numerically, and I find $f(0.124) = 6.61$. Since $f(T/T_c)$ is a uniformly decreasing function of T/T_c , I obtain (using (12)) the condition for no nucleation until the present era

$$260 \leq F_c/T \leq (1.22)(6.61) \lambda^{-3/4} , \tag{35}$$

giving a bound on λ ,

$$\lambda \leq 0.010, \text{ or } -0.010 \leq \bar{\lambda}(0) \leq 0 . \tag{36}$$

Fine Tuning?

Does the bound (36) represent a substantial fine tuning? In an attempt to answer this question, I have presented in Figure 1 the renormalization group flow in the $\bar{\lambda} - \bar{\lambda}_2$ plane over less than or equal to a decade $[(t \leq \ln(10)) \rightarrow (t = 0)]$ for those trajectories which begin in the stability region and satisfy the requirement (36). It is seen that (1) a reasonable piece of the coupling constant phase space is available (2) the couplings are perturbative but not particularly small over much of the RG flow and (3) over some of the phase space $\bar{\lambda}$ reaches values of the order of the electroweak Higgs coupling. (In the same normalization, $\lambda_H = 0.08 (m_H/100 \text{ GeV})^2$.)

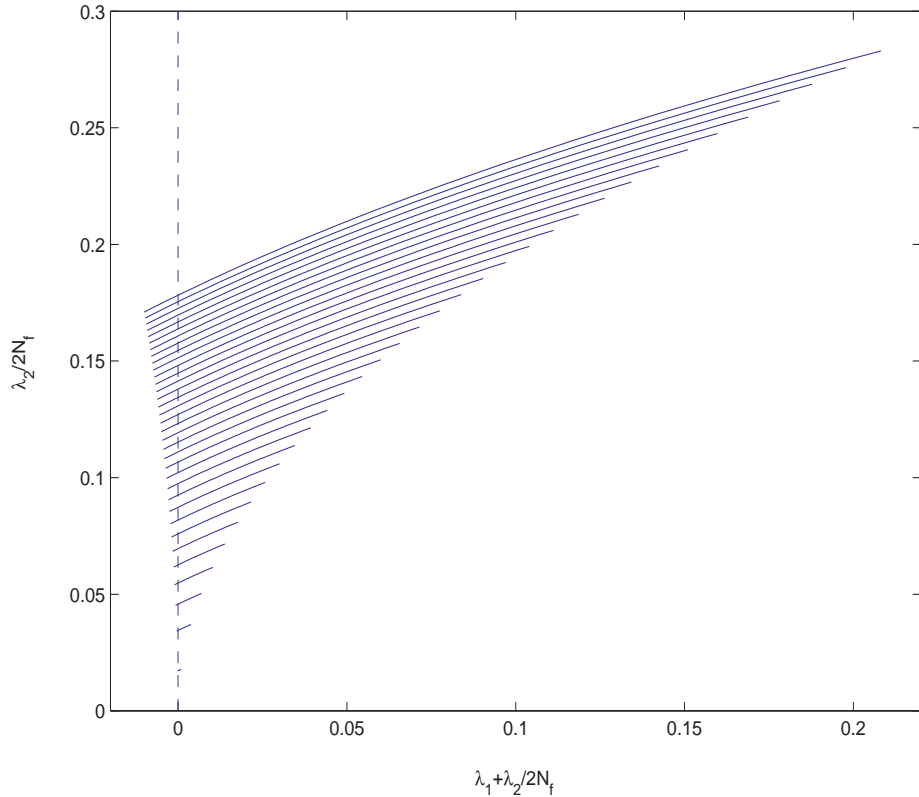


Figure 1: Basin of attraction (for $t \leq \ln(10)$) to relevant parameter space (Eq. (36)) at $t = 0$. Region to right of dashed line is stable for all field directions.

A summary of results and concluding remarks follows.

4 Discussion and Conclusions

(1) A model has been presented which can generate a cosmological constant of magnitude to account for $\Omega_\Lambda \simeq 0.7$ during the present era. The present quasi-deSitter phase is driven by the false vacuum energy associated with the supercooled phase of an incomplete chiral phase transition in a hidden gauge theory. The very small energy scale $\rho_\Lambda^{1/4} \simeq 2 \times 10^{-3} \text{ eV}$ for the vacuum energy

appears as the strong interaction scale for a hidden $SU(2)'$ whose coupling runs from GUT energies, with coupling at GUT not quite unifying with the standard model couplings. Having this $SU(2)'$ and the standard model fields reside on different branes presents a simple solution to this disparity.

(2) An unchanging vacuum energy has some advantages over evolving primordial scalar fields as an origin of the present near-deSitter phase: the problem of protecting the tiny curvature of the potential [45, 46] is circumvented, as is the necessity (in quintessence models) to control the contribution of dark energy during nucleosynthesis [13].

(3) BBN considerations imply large supercooling in the metastable quark-gluon plasma of the present phase. Although present lattice simulations indicate only small supercooling in the *quenched* approximation of Yang-Mills theory, theoretical considerations indicate that large supercooling is possible in the presence of chiral condensates. A linear sigma model with two quartic couplings was analyzed in which all symmetry breaking takes place via dimensional transmutation. In this model the existence of substantial supercooling does not require fine tuning in the coupling constant space. The existence of a large interface energy in an $SU(2)$ theory with $N_f = 4$ dynamical quarks would provide incisive support for this model.

(4) A pivotal requirement in this scenario is the near-scale invariance of the chiral lagrangian. This was briefly discussed in the text in terms of the ratio N_f/N ($= 2$ in the present work). A hint of the continuous approach to a conformal phase may be suggested in the observed weakening of the chiral phase transition in a lattice study [41] of four flavor QCD with N_f/N as small as $4/3$.

(5) The bound $\lambda \leq 0.010$ (Eq. (36)) does not change drastically with a tighter bound on the extra effective number of neutrinos. For example, with a requirement $\Delta N_{eff} \leq 0.1$, I find instead of (36) the bound $\lambda \leq 0.0086$. This would also involve a small amount of extra reheating in the visible sector between the GUT and electroweak scales.

(5) Although the phase transition discussed in this paper is long overdue, it may not be catastrophic when it occurs. The nucleation occurs via very thick-walled bubbles, so that the drastic shock-wave scenario depicted in the thin-walled case [43] is perhaps not inevitable.

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